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D-Matter

Gary Shiu and Lian-Tao Wang

Department of Physics, University of Wisconsin, Madison, WI 53706, USA

Abstract

We study the properties and phenomenology of particle-like states originating from D-branes whose spatial dimensions are all compactified. They are non-perturbative states in string theory and we refer to them as *D-matter*. In contrast to other non-perturbative objects such as 't Hooft-Polyakov monopoles, D-matter states could have perturbative couplings among themselves and with ordinary matter. The lightest D-particle (LDP) could be stable because it is the lightest state carrying certain (integer or discrete) quantum numbers. Depending on the string scale, they could be cold dark matter candidates with properties similar to that of wimps or wimpzillas. The spectrum of excited states of D-matter exhibits an interesting pattern which could be distinguished from that of Kaluza-Klein modes, winding states, and string resonances. We speculate about possible signatures of D-matter from ultra-high energy cosmic rays and colliders.

Email addresses: shiu@physics.wisc.edu, liantaow@pheno.physics.wisc.edu.

I. INTRODUCTION

In addition to their pivotal role in elucidating non-perturbative aspects of string theory, in recent years D-branes have become ubiquitous in particle phenomenology and cosmology. In the brane world scenario [1, 2, 3, 4, 5, 6], the Standard Model gauge bosons and matter fields are localized on some space-filling branes whereas only gravity and closed string modes (such as the radions) can propagate in the bulk, thus offering the possibility of a lower fundamental energy scale. It is worth noting that in this framework the D-branes are part of the string theory vacuum. The Standard Model particles and the bulk modes are excitations above the vacuum described by *perturbative* string states in the background of D-branes. Therefore, even though the branes themselves are non-perturbative objects, the physics of the brane world scenario can be studied using standard techniques of string perturbation theory.

In this paper, we explore yet a different phenomenological feature of D-branes. Unlike the background D-branes in the context of brane world, the D-branes we consider are not space-filling but rather behave as point particles in four dimensions. These particle-like states arise when all spatial dimensions of Dp-branes (irrespective of p) are wrapped around the compact space. We refer to them collectively as *D-matter*. As we will discuss in more detail, these D-matter states appear quite naturally in four-dimensional string models. They are non-perturbative objects with mass $m \sim M_s/g_s$ where g_s is the string coupling and so at weak coupling they are heavier than the perturbative string states. However, in order to produce the observed gauge and gravitational couplings, g_s is not arbitrarily small but is typically of order one. Hence, the D-matter states can be sufficiently light to be phenomenologically and cosmologically relevant. The D-matter states are dynamical degrees of freedom and so, in contrast to the background D-branes, we will treat them as excitations above the vacuum.

The stability of D-branes is due to the charges that they carry. Depending on their types (BPS or non-BPS), D-branes could carry integral or torsion (discrete) charges. In fact, it is because of the torsion charges that the complete spectrum of D-branes should be classified by K-theory [7, 8, 9]. Among the D-matter states, the lightest D-particle (LDP) is stable because it is the lightest state carrying its specific charges. Therefore, just like the lightest supersymmetric partners (LSPs) in supersymmetric models, the LDPs are possible candidates for cold dark

matter. As we will show, an important difference between the D-matter states and other non-perturbative objects (such as magnetic monopoles [61]) is that they could have *perturbative* couplings. Hence, the LDPs are weakly interacting and so they could be candidates for wimps or wimpzillas depending on the string scale. We also comment on the case where the D-matter states are unstable (i.e., the spectrum of stable D-branes does not contain four-dimensional particle-like states). Although they are not long-lived to be cosmologically relevant, they could be produced at colliders and give rise to interesting signatures.

In addition to D-matter which are particle-like states (from a four-dimensional perspective), other types of stable defects (such as cosmic strings and domain walls) could also exist in a string vacuum. Given a specific string model, it is straightforward to deduce the complete spectrum of such defects. In this paper, we focus on the general properties of D-matter and point out their relevance to phenomenology. We also stress that the cosmological constraints on the various kinds of stable defects could provide important additional criteria on the viability of a string model (and hence in estimating the number of realistic string vacua [11]). For example, we can rule out a string vacuum if there exist stable electrically charged D-matter states.

This paper is organized as follows. In Section II, we describe the properties of D-matter including its stability, mass, and interactions. In Section III, we discuss the phenomenological implications and some possible signatures of D-matter. Finally, we end with some discussions and conclusion in Section IV. Some details of the four-point heterotic string amplitude which is dual to the annihilation cross-section of D-matter states are relegated to the appendix.

II. PROPERTIES OF D-MATTER

A. Mass of D-matter

D-matter has a very simple origin. A Dp-brane sweeps out a $p+1$ dimensional worldvolume. From a four-dimensional perspective, a D-brane with all of its spatial dimensions compactified will sweep out a worldline and behave like a point particle.

In this section, we will consider the constraints on the mass of D-matter. Since D-matter states arise from Dp-branes wrapping around p compactified dimensions, their masses are given

by:

$$m_D = \frac{M_s^{p+1} V_p}{g_s} \ , \tag{1}$$

where V_p is the volume of the p-cycle (i.e., p-dimensional subspace of the internal manifold) that the Dp-branes wrapped around. The D-matter states are non-perturbative objects whose masses are inversely proportional to g_s. At weak coupling, they are heavier than the perturbative string states. However, g_s cannot be arbitrarily small for otherwise the Planck mass M_P is not finite. This can be seen from the relation:

$$M_P^2 = \frac{M_s^8 V_6}{g_s^2} \tag{2}$$

where V_6 is the overall volume of the internal manifold. Hence, m_D is not much heavier than M_P. As we shall see shortly, if the dimensions transverse to the D_p brane have size $R_\perp$ larger than the string scale $M_s R_\perp >> 1$, the resulting D-matter state can be much lighter than M_P.

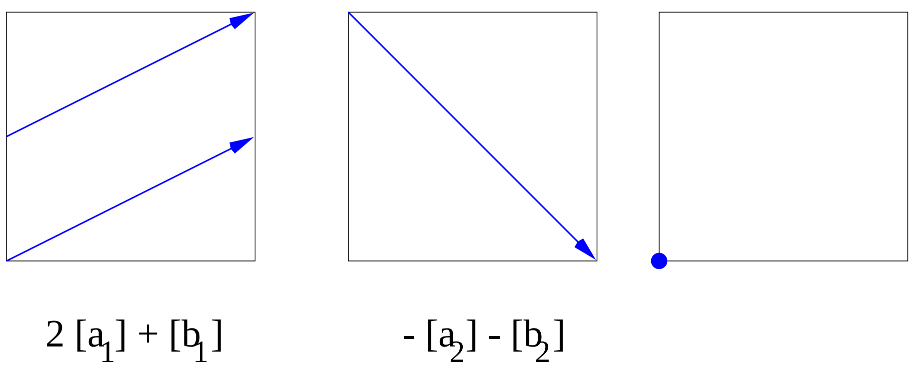

FIG. 1: A stable D2-brane wrapping around the 2-cycle $(2[a_1] + [b_1]) \times (-[a_2] - [b_2])$.

We can obtain a more quantitative estimate of m_D. Although in general there is no direct relation between V_p and V_6, for simple compactifications such as toroidal and orbifold backgrounds, we can show that generically $M_s/g_s \lesssim m_D \lesssim M_P$. For illustrative purpose, let us consider the internal manifold to be a product of tori $T^2 \times T^2 \times T^2$ (or its orbifold thereof) with sizes R_i where $i = 1, \ldots, 3$, then

$$M_P^2 = \frac{M_s^8 \prod_{i=1}^3 R_i^2}{g_s^2} \ , \tag{3}$$

$$m_D = \frac{M_s^{p+1} V_p}{g_s} \sim \frac{M_s^{p+1} \prod_{j=1}^p L_j}{g_s}, \tag{4}$$

where we assume that the p-cycle that the Dp-brane wrapped around is factorisable into a product of one-cycles with lengths L_j (which are related to R_i by the wrapping numbers). Let us denote a canonical basis of one-cycles in each T^2 by $[a_i]$ and $[b_i]$ (where $i = 1, 2, 3$) respectively. As an example, suppose that the D-matter state originates from a stable D2-brane wrapped around a two-cycle $(n_1[a_1] + m_1[b_1]) \times (n_2[a_2] + m_2[b_2])$ as shown in Fig. 2, then $p = 2$, and $L_i = \sqrt{n_i^2 + m_i^2} R_i$. The wrapping numbers are typically of order one and so $L_i \sim R_i$. Note that D-brane has a size of the order of M_s^{-1} and so if $R_i < M_s^{-1}$, it is more appropriate to go to the T-dual picture:

$$R_i \to \frac{1}{M_s^2 R_i} \qquad g_s \to \frac{g_s}{M_s R_i} \tag{5}$$

and the Dp brane becomes a $D(p-1)$ or a $D(p+1)$ brane depending on whether R_i is along or transverse to the worldvolume of the Dp brane. Therefore, there is a natural lower bound for the size of the compact dimension $R \geq M_s^{-1}$. As a result, the mass of a D-matter state is bounded below by M_s/g_s. For the same reason, $m_D \lesssim M_P$ and furthermore if some dimensions transverse to the Dp branes are large, the D-matter state can be much lighter than M_P.

The mass of D-matter depends on the value of M_s. In the brane world scenario, M_s is not tied to M_P and since $g_s \sim O(1)$, the D-matter mass can be anywhere between the TeV scale and the Planck scale. As we will discuss, the D-matter states could have some interesting phenomenological consequences if they are sufficiently light.

B. Stability of D-matter

Let us denote a D-brane which has r Neumann directions in our usual four-dimensional spacetime and s Neumann directions in the compact space as a $D(r, s)$-brane. In this notation, there exist stable D-matter states if the spectrum of stable D-branes contains some $D(r, s)$-branes with $r = 0$ (irrespective of s).

The complete spectrum of stable D-branes in a given string theory background is rather rich and in general non-trivial to work out. This is because contrary to naive expectations, D-brane charges are not classified by cohomology but instead by K-theory [7, 8, 9]. In addition to the BPS D-branes, there are generically non-BPS D-branes that are stable [8, 12, 13, 14]. The existence of such stable non-BPS branes is further exemplified by some concrete orbifold

[15, 16] and orientifold [17, 18] models.

The spectrum of stable D-branes can be divided into the following types:

- *BPS states* are stable because they carry Ramond-Ramond (RR) charges, which are conserved charges associated with gauge fields coming from the RR sectors of string theory. The RR charges are integrally quantized, analogous to the winding number of a monopole solution. As a result, the BPS state with the smallest unit of charge ($n = 1$) is stable. An example of a stable BPS brane is the $D0$-brane in Type IIA string theory (which behaves as particles even before compactification).

- *Stable non-BPS states* which can be further divided into the following types:

  1. Non-BPS branes which carry twisted RR charges, i.e., they are charged under the RR gauge fields localized at orbifold fixed points. These twisted RR charges are quantized, just like the untwisted RR charges carried by BPS branes. Therefore, the lightest state carrying some given twisted RR charges is stable.

  2. Torsion charged D-branes. They are non-BPS branes which are stable even though they do not carry any RR charges. Rather, they carry some torsion (discrete) charges. It is precisely these torsion charges that K-theory differs from cohomology. An example of such torsion-charged D-branes is the stable D-particle in Type I string theory [12]. Although $D0$-brane is unstable in Type IIB string theory, the tachyon which signals the instability is projected out by the orientifold projection. Hence, in Type I string theory (which can be obtained from Type IIB theory by an orientifold projection), the D-particle is stable and in fact carries a $\mathbf{Z}_2$ torsion charge.

The lightest D-matter states arising from some stable D-branes wrapping around the compactified dimensions are stable because they are the lightest states carrying certain (integer or discrete) quantum numbers. In what follows, we will refer to them as LDPs (lightest D-particles).

C. Interactions of D-matter

Let us discuss how the D-matter states interact with each other and with the Standard Model particles. To be concrete, we consider a specific case where the D-matter states are the stable D-particles in Type I string theory and the Standard Model is embedded on a set of D9-branes. Although we focus our analysis on this particular case for illustrative purposes, we expect the properties of their interactions discussed below are applicable to a general class of D-matter states. The general rules for computing amplitudes involving the stable Type I D-particles have been given in [8, 19] and such rules have been applied to studying their gauge and gravitational interactions in [20]. In the brane world scenario, the gauge bosons and the matter fields of the Standard Model can in general be localized differently in the extra dimensions. For example, in the context of intersecting brane world [21, 22, 23, 24, 25, 26, 27, 28, 29, 30, 31, 32, 33, 34], the gauge bosons are localized on several different stacks of Dp'-branes with $p' \geq 3$ whereas chiral matter fields are localized at the brane intersections. For simplicity, however, let us consider the case where all Standard Model fields are localized on a single stack of Dp'-branes, keeping in mind that chiral fermions can still arise in this scenario when the D-branes are located at some singularities of the internal manifold.

The worldvolume of the Dp'-branes (where $p' = 9$ in this particular case, in general $p' \geq 3$) contains our usual four-dimensional space-time and in general also some compactified dimensions. The D-matter state propagates in the background of the higher-dimensional Dp'-branes where the Standard Model is localized. It is worth emphasizing that in this "branes-within-branes" picture, the D-matter states are not part of the background but instead its *excitations*. Since g_s is non-zero, their masses are finite and so they are dynamical degrees of freedom propagating in some background D-branes.

In this simple setup, the Standard Model fields are open string excitations with both end-points of the open strings attached to the same Dp'-brane. These open string states are denoted by $C_{p'p'}$. There are also open strings that stretched between the propagating D-matter state and the background Dp'-brane and we denote these states by $C_{0p'}$. As illustrated in Fig. 2, the $C_{0p'}$ states couple to the Standard Model fields in the $C_{p'p'}$ sector. Hence, there is an effective coupling of D-matter to the gauge bosons (and other Standard Model fields) via the

interactions between perturbative open string states. The D-matter states are charged under the gauge groups localized on the Dp'-branes. Therefore, they couple with matter fields on the branes through gauge interactions.

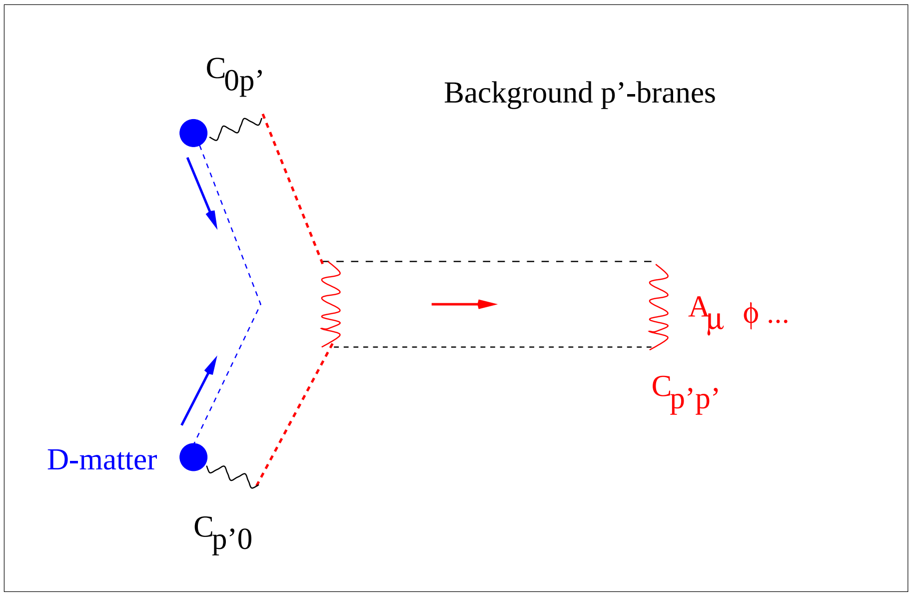

FIG. 2: D-matter can interact via perturbative string states.

Such D-matter states have gauge couplings and so they can interact with each other by the exchange of gauge bosons. An important aspect of the gauge (or matter) coupling is that it is perturbative. At lowest order, the interaction of such D-matter with the Standard Model fields is given by the disk diagram in Fig. 3 where $V_{0p'}$, $V_{p'0}$ and $V_{p'p'}$ are vertex operators for open strings in the $0p'$, $p'0$, and $p'p'$ sectors respectively. The amplitude of this diagram is proportional to $g_{YM} \propto g_s^{1/2}$ [62]. Higher order couplings between D-matter and the Standard Model fields are suppressed by the higher power of g_s. Therefore, from power counting, the S-matrix describing the interactions of D-matter has an expansion in terms of g_s.

However, we should be more careful in writing down the effective coupling. From power counting, the amplitude describing the interaction of D-matter with Standard Model fields (whose lowest order contribution is shown in Fig. 3) has an expansion

$$g_s^{\frac{1}{2}} F(s,t,\alpha') + g_s G(s,t,\alpha') + \ldots \tag{6}$$

We see that the effective coupling extracted from this expression is $(g_s^{\frac{1}{2}}F(s,t,\alpha')+...)\bar{\Psi}A^\mu\Psi$. In other words, there is a "form factor" which effectively "renormalizes" the coupling constant. In some kinematical region, such as forward elastic scattering, we can safely ignore this correction due to the low momentum exchange. In fact, in this case, the deviation from a lowest order field theory result is expected to be suppressed by powers of q/M_S where q is the momentum exchange. However, in some other kinematical region of interest, such as s-channel annihilation, the correction is not expected to be small due to the fact that the energy involved exceeds the string scale. In this regime, we expect corrections of stringy nature to become significant. The question is then the magnitude of such corrections.

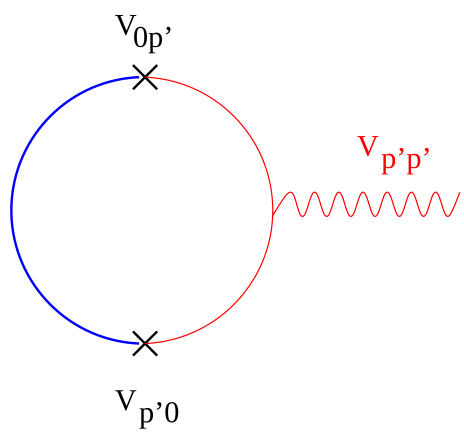

FIG. 3: At lowest order, the interaction of D-matter with the Standard Model fields is described by a disk diagram with the insertion of open string vertex operators.

The techniques for calculating the full annihilation cross section (i.e., including the stringy corrections) of D-branes have not been fully developed. Nevertheless, we can estimate the size of the stringy corrections (i.e., α' effects) by studying this annihilation process in the dual heterotic string picture. According to the Type I - heterotic duality [35], the stable D-particle in Type I string theory (an example of D-matter) is dual to a massive heterotic string state [12]. In the dual heterotic string picture, the annihilation process of interest involves only perturbative string states and so the corresponding amplitude can be calculated using standard techniques

	't Hooft-Polyakov Monopole	D-Matter
Mass	$\frac{<\phi>}{g_{YM}} \sim \frac{M_X}{g_{YM}^2}$	$\frac{M_s}{g_s} \sim \frac{M_s}{g_{YM}^2}$
$(\text{size})^{-1}$	$\lambda < \phi >$ $g_{YM} < \phi >$	$g_s^{\alpha} M_s$ $\alpha = \begin{cases} -1/3 & \text{brane-probe} \\ 0 & \text{string-probe} \end{cases}$
Interaction	$\propto \mu_m = \frac{n}{g_{YM}}$ where $n \in \mathbf{Z}$	$\propto g_{YM}$

TABLE I: Comparison between 't Hooft-Polyakov monopole and the D-Matter discussed in this section. Here, $< \phi >$ denotes the vacuum expectation value of the scalar field ϕ of the monopole configuration, λ is its coupling constant, and M_X is the symmetry breaking scale. The size of D-matter depends on the probe since D-branes can probe smaller distances than fundamental strings [36, 37].

of string perturbation theory. The details of this heterotic string amplitude calculation are given in the Appendix. The string amplitude can then be compared with the field theory result (which corresponds to $\alpha' \to 0$). As shown in Fig. 5, the α' corrections over a wide range of energies are of the same order as the field theory result and so they cannot be neglected.

Hence, we expect that the D-matter gauge coupling would have an $\mathcal{O}(1)$ correction from the form factor. This correction should be small when we consider elastic scattering with small momentum exchange (small comparing to string scale). The correction would be important when we consider s-channel annihilations. However, we do not expect order of magnitude enhancement or suppression. After all, the energy scale involved is only slightly above the string scale and so the corrections only come from the lowest few string resonances.

It is interesting to compare the properties of such D-matter with that of a 't Hooft-Polyakov

monopole in spontaneously broken gauge theory which is also a non-perturbative object. They have similar properties as far as their masses and sizes go. However, the 't Hooft Polyakov monopole carries a magnetic charge $\mu_m \propto 1/g_{YM}$ under the unbroken $U(1)$. Therefore, unlike the D-matter considered here, the long range gauge interactions between 't Hooft Polyakov monopoles are non-perturbative [63]. A comparison between 't Hooft Polyakov monopole and the D-matter discussed here is summarized in Table I.

III. PHENOMENOLOGY OF D-MATTER

In this section, we study some phenomenological features of D-matter. The spectrum and quantum numbers of the D-matter states are model-dependent. Instead of studying their properties in the context of a specific model, we focus on the following two generic features.

A. Cold Dark Matter Candidate

As discussed above, the LDPs are stable and so they could be candidates for the cold dark matter of our universe. Their properties as dark matter candidates depend on the value of the string scale. In the conventional high string scale scenario ($M_S \sim 10^{16}$ GeV), the LDPs have masses greater than 10^{16} GeV. They are presumably removed by inflation (just as the GUT monopoles) and would not be produced significantly during reheating. Therefore, we consider the following two scenarios with intermediate and low string scales, respectively.

1. *Intermediate string scale:* $M_s \sim 10^{11} - 10^{12}\ GeV$

   In this scenario, the mass of the LDP will be of the order of $10^{11} - 10^{12}$ GeV. It therefore fits in the category of wimpzillas, i.e., super-massive wimps. Since their masses are higher than the reheating temperature, a key question is whether they can be produced in significant amount to account for the cold dark matter in the universe. This question has been studied in a series of papers [38, 39, 40, 41, 42, 43, 44]. Several scenarios have been proposed to produce them including preheating , non-adiabatic expansion of the universe, and normal reheating. Each of these mechanisms has its own special properties. It is interesting that sufficient number of the superheavy LDPs can be produced through

normal reheating. From Ref. [42], the relic abundance of such a state produced during reheating is

$$\Omega_D h^2 = M_D^2 < \sigma v > \left(\frac{g_*}{200}\right)^{-3/2} \left(\frac{2000 T_{RH}}{M_D}\right)^7 . \tag{7}$$

The absence of the naively expected $\exp(-M_D/T_{RH})$ suppression is due to the fact that during the reheating process, the maximum temperature is much higher than the reheating temperature. Using the expression for the production cross section, we see that states as heavy as $10^3 T_{RH}$ could be produced in an appropriate amount to account for the cold dark matter in the universe.

On the other hand, if the LDPs only couple to the light states (or inflaton) with much weaker (or not at all) couplings (such as gravitational strength couplings), they could be produced in a "non-adiabatic" expansion stage at the end of inflation [39].

An interesting consequence of super-heavy dark matter is that their annihilation could give rise to ultra-high energy cosmic rays (UHECRs) with energies $> 10^{20}$ eV. The possibility of decaying super heavy dark matter within galactic halo as a source of UHECR events[45, 46] which exceed the GZK cutoff [47, 48] has been studied [49, 50]. Although the rate of UHECR from the annihilation of D-matter is expected to be smaller than the measured value (without special assumptions about their couplings), it is possible that they may account for some of the events provided that some special features of the local density of dark matter are satisfied. If this is indeed the source of the UHECRs, another interesting correlated signal will be an enhanced possibility of detecting high energy cosmic ray neutrinos in Amanda II and IceCube [51].

2. *Low string scale: $M_s \sim$ TeV*

   Such a low string scale implies that our stable D-matter has a mass around TeV as well. Suppose they interact with each other and the light degree of freedoms in the thermal bath with the strength of gauge interactions, then they have many similarities with the well studied thermal relics (such as the LSP in MSSM). They will be produced and in thermal equilibrium with the thermal bath during reheating. Then as the universe expands, they will follow a standard freeze out procedure. A key quantity which determines the relic

abundance is their annihilation cross section $< \sigma v >_A$. Since the D-matter behaves as a particle in 4-dimensions, we could write down an effective field theory for its coupling to the gauge bosons. In particular, there will be a term of the form $g_D \bar{D} D A$ (we suppress the Lorentz structure which may contain gamma matrix or spacetime derivative). The coupling constant g_D should be determined by a string theoretic calculation (generically, it has the strength of a normal gauge coupling). The thermal relic abundance as a result of the freeze out process is

$$\Omega_D h^2 \sim \frac{3 \times 10^{-6}}{\alpha_D^2} \left(\frac{M_D}{100\mathtt{GeV}} \right)^2 . \tag{8}$$

From this we see that for a gauge coupling strength interaction, we can have a relic abundance appropriate for the CDM. Notice that this case is very similar to the well studied LSP cold dark matter in the context of MSSM.

Let us discuss the prospects of detecting the LDPs. Direction detection of dark matter[52] relies on the elastic scattering of dark matter off the nuclei within the detector. As discussed in Section II C, we do not expect a sizable modification of the effective gauge coupling in this channel. Therefore, we should be able to calculate their rate reliably. However, there is a subtlety we need to keep in mind. The dark matter-nucleon cross-section is sensitive to the spin of the dark matter particle [52]. For example, in MSSM, a neutralino (which is a Majorana fermion) and a sneutrino (which is a boson) have very different direct detection cross-sections. Similarly, D-matter states which are bosonic and fermionic respectively [64] will have different dark-matter-nucleon cross-sections.

Indirect detection of the dark matter relies on detecting the cosmic rays resulting from their annihilation within the galactic halo. As discussed in Section II C, we expect the annihilation cross section to be modified with respect to the field theory result. This may have observable effects in high energy cosmic rays such as an excess (or deficiency) of cosmic ray flux.

B. Excited States and Mass Splittings

Another interesting feature of D-matter is the mass spectrum of its excited states. A tower of excited states can be constructed by dressing the D-matter with excitations of open strings attached to it. These D-matter excited states could be a charged state or a singlet under the gauge symmetries of the background D-branes depending on whether one or both endpoints of the open strings are attached to the D-matter [65]. Thus, these two types of excited states could lead to different patterns of decay modes. An example of such an excited state is shown in Fig. 2. Schematically, an excited D-matter state can be written as

$$|D_0 > \times |\mathcal{V} >, \tag{9}$$

where $|D_0 >$ is the "bare" D-matter and $\mathcal{V}$ is the vertex operator of some excited open string state. The mass of such an excited state is given by [66]:

$$M_{D^*}^2 = M_{D_0}^2 + nM_s^2, \tag{10}$$

where $M_{D_0} \sim M_s/g_s$. This is because a string resonance contributes to the *square* of the mass of a state by an amount of nM_s^2 where n is the resonance level.

Let us compare this tower of excited states with the familiar Kaluza-Klein (KK) modes, winding modes, and string resonances:

$$\begin{aligned} M_{KK}^2 &= g_s M_s^2 + n^2 M_C^2, \\ M_{\text{winding}}^2 &= g_s M_s^2 + n^2 \frac{M_s^4}{M_C^2} \\ M_{\text{resonances}}^2 &= g_s M_s^2 + nM_s^2 \\ M_{D^*}^2 &= \frac{M_s^2}{g_s^2} + nM_s^2, \end{aligned} \tag{11}$$

where M_C is the compactification scale. The zero point masses of the KK modes, winding modes, and string resonances are taken to be $g_s M_s^2 \sim g_{YM}^2 v^2$, which presumably come from some spontaneous symmetry breaking. A sample spectrum of the various towers of excited states is shown in Fig. 4. In the case of a low string scale scenario, these different towers of new particles (whose mass spectra have distinct patterns) can have interesting observable signatures at colliders. Note that the open string states attached to the D-matter can in general carry

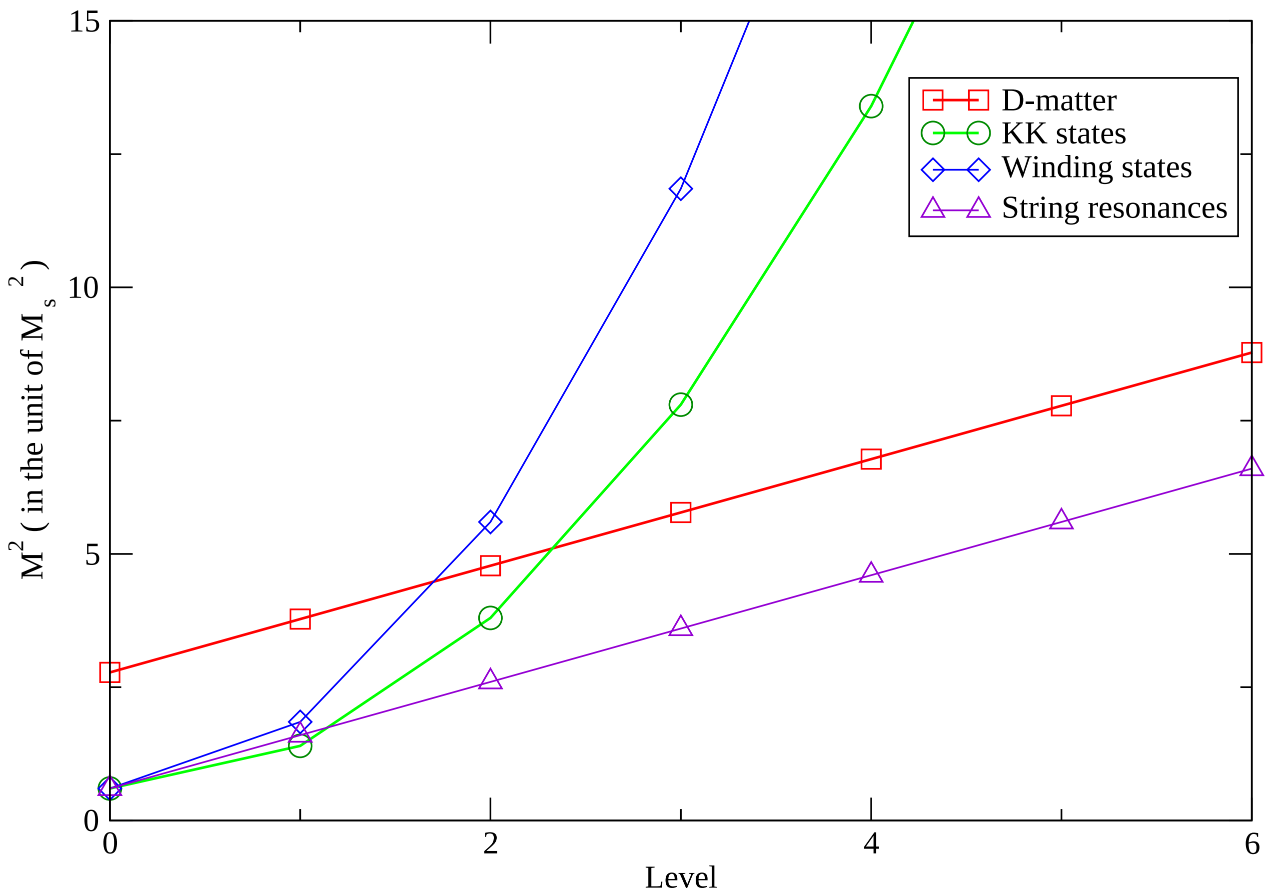

FIG. 4: Different towers of excited states. The masses are given in unit of M_S. For illustrative purpose, we choose a representative set of parameters: the compactification scale $M_C \sim R^{-1}$ is taken to be $0.8M_s$, and the string coupling $g_s \sim 0.6$. With some symmetry breaking scenario in mind, we choose a zero point energy for the zero modes to be $g_s M_s$.

momentum and/or winding quantum numbers as well, but the details depend on the origin of the D-matter and how the Standard Model is embedded [67]. For illustrative purpose, we consider the case where the D-matter is a stable D-particle (e.g., in Type I string theory) and the Standard Model is embedded on a set of D9-branes. In this case, the open string states

have no momentum or winding numbers. In general, the possibility of momentum/winding quantum numbers will give rise to an even richer spectrum than the one shown in Fig. 4.

For simplicity, we have taken the zero point mass of the D-matter state to be $m_D \sim M_s/g_s$. This is based on the crude assumption that the volume of the cycle V_p that the stable D-brane wraps around is of the order of M_s^{-1}. In general, $m_D \sim M_s^{p+1} V_p / g_s$. Therefore, the zero point mass of D-matter carries some useful information about the compact dimensions, which may allow us to distinguish between isospectral manifolds (different manifolds which nonetheless have the same spectrum) that are indiscernible by the KK and winding modes [53].

IV. DISCUSSIONS AND CONCLUSION

In this paper, we study the phenomenology of D-matter, i.e., particle-like states originated from D-branes whose spatial dimensions are all wrapped around the compact space. An interesting feature of D-matter is that although they are non-perturbative objects, they could have perturbative couplings with each other and with the Standard Model fields. Therefore, the D-matter states are weakly interacting and the lightest stable D-matter states (whose stability is due to integral or torsion charges they carry) are candidates for wimps or wimpzillas depending on the string scale. In the case of a low string scale, there are also potentially rich collider phenomenologies. Exploring the mass pattern of the lower lying D-matter states could reveal additional information about the structure of the compact dimensions.For illustrative purposes, we consider the simple case where all the Standard Model particles are localized on a single set of Dp'-branes. It would be interesting to explore the phenomenology of D-matter in more realistic brane world models such as models with background intersecting D-branes [21, 22, 23, 24, 25, 26, 27, 28, 29, 30, 31]. In these models, the gauge bosons and matter fields are localized differently in the extra dimensions (e.g., in different codimensions) and hence they will couple with different strengths to the D-matter.

Another interesting direction motivated by this work is to compute the spectra of stable D-branes (and their corresponding K-theories) in some realistic brane world models, e.g., certain four-dimensional $N = 1$ supersymmetric orientifold models whose perturbative string spectra contain the Standard Model or Grand Unified Theories [21, 22, 23, 24]. The K-theories which

classify the D-brane spectra have been computed for some concrete orientifold models, e.g., orientifolds with $N = 2$ supersymmetries (in a four-dimensional sense) in [17, 18]. It would be interesting to extend these analyses to models with $N = 1$ supersymmetry where chiral fermions can appear. Such analyses could also help to better understand the types of defects (from a four-dimensional perspective) that are formed toward the end of brane inflation [54], besides the cosmic strings that are generically produced [55, 56, 57, 58, 59]. Therefore, in addition to the conditions one often imposes (such as appropriate gauge groups, three generations of chiral matter, etc) in estimating the statistics of realistic string theory vacua [11], the cosmological constraints on the existence of stable defects (such as cosmic strings and domain walls) may further restrict the space of viable string models.

APPENDIX A: FOUR-POINT AMPLITUDES IN HETEROTIC STRING DUAL

The stable non-BPS D-particles in Type I are dual to the first massive string states of the $SO(32)$ heterotic string [12]. The vertex operators of such massive string states in the R-sector (spacetime fermions) in the -1/2 picture are given by:

$$V^{-1/2}(z,\overline{z}) = e^{-\phi(z)/2} e_\mu u_\alpha i\partial_z X^\mu \theta^\alpha(z) e^{ik\cdot X(z,\overline{z})} C_K e^{i\frac{K}{\sqrt{2\alpha}}\cdot\tilde{X}(\overline{z})} \tag{1}$$

and

$$V^{-1/2}(z,\overline{z}) = e^{-\phi(z)/2} e_\mu u_\alpha i\partial_z \psi^\mu \theta'^\alpha(z) e^{ik\cdot X(z,\overline{z})} C_K e^{i\frac{K}{\sqrt{2\alpha}}\cdot\tilde{X}(\overline{z})} \tag{2}$$

where $k^2 = -M^2 = -4/\alpha'$ and $K^2 = 4$. Here θ_α, θ'_α are the ten-dimensional spin fields with opposite chirality (because of the GSO projection and the fact that ψ changes the fermion number by 1), e_μ is the polarization (the physical state condition gives $e \cdot k = 0$), $X^\mu(z,\overline{z})$ for $\mu = 0, \dots, 9$ are the 10-dimensional coordinates for both left- and right-movers, $\tilde{X}^I(\overline{z})$ for $I = 1, \dots, 16$ are the 16-dimensional internal coordinates corresponding to the $SO(32)$ lattice, C_K is a cocycle factor, and ϕ is the boson arising from the bosonization of the superconformal ghosts. Using the fact that the conformal dimension of $e^{q\phi}$ is $h = -q(q+2)/2$, it is easy to see that the above vertex operators have conformal dimensions $(h,\overline{h}) = (1,1)$. We take $K = (\pm\frac{1}{2}, \pm\frac{1}{2}, \dots, \pm\frac{1}{2})$ with an even number of + sign which clearly gives $K^2 = 4$. These massive states indeed transform in the spinor rep. of $SO(32)$. In addition to the $SO(32)$

quantum numbers, the two types of vertex operators above together contribute to the **128** representation of the Lorentz group. For simplicity, we focus on the first type of vertex operators in Eqn.(1) in the calculation of four-point amplitudes.

The scattering of two D0-branes to two fundamental open string states corresponds to the scattering of two massive string states to two massless states in the heterotic dual. The tree-level diagram has ϕ-ghost number -2, and so the four-point amplitude can be computed with all external states in the -1/2 picture. In other words, the scattering amplitude of two massive spacetime fermions in the spinor reps of SO(32) into two gauginos:

$$A_4 \sim \mathcal{C}_0 \hat{N}^4 < c\overline{c} V_1^{-1/2}(z_1,\overline{z}_1) V_2^{-1/2}(z_2,\overline{z}_2) V_3^{-1/2}(z_3,\overline{z}_3) V_4^{-1/2}(z_4,\overline{z}_4) > \tag{3}$$

where $\mathcal{C}_0$ and $\hat{N}$ are normalization factors defined in [20], V_1 and V_2 are the vertex operators for the massive string states in the spinor rep. as in Eqn. (1) above, V_3 and V_4 are vertex operators for the gauginos as follows:

$$V^{-1/2}(z,\overline{z}) = e^{-\phi(z)/2} u_\alpha \theta^\alpha(z) e^{ik\cdot X(z,\overline{z})} \mathcal{O}(\overline{z}) \tag{4}$$

with $k^2 = 0$. $\mathcal{O}(\overline{z})$ is the operator for the $SO(32)$ degrees of freedom, which is equal to $\overline{\partial}\tilde{X}^I(\overline{z})$ for the Cartan generators and $C_K e^{i\frac{K}{\sqrt{2\alpha}}\tilde{X}(\overline{z})}$ with $K^2 = 2$ for the remaining 480 generators. We can take K of the form $(0,\ldots,\pm1,0,\ldots,0,\pm1,0,\ldots,0)$. In either case, $\overline{h}_{\mathcal{O}} = 1$, so it is easy to see that $(h,\overline{h}) = (1,1)$.

Let $\mathcal{V}_\alpha(z) = e^{-\phi/2}\theta_\alpha(z)$, the relevant correlation function (see, e.g., eqn (12.4.19) of [60]) is:

$$< \mathcal{V}_\alpha(z_1)\mathcal{V}_\beta(z_2)\mathcal{V}_\gamma(z_3)\mathcal{V}_\delta(z_4) >= \frac{(C\Gamma^\mu)_{\alpha\beta}(C\Gamma_\mu)_{\gamma\delta}}{2z_{12}z_{23}z_{24}z_{34}} + \frac{(C\Gamma^\mu)_{\alpha\gamma}(C\Gamma_\mu)_{\delta\beta}}{2z_{13}z_{34}z_{32}z_{42}} + \frac{(C\Gamma^\mu)_{\alpha\delta}(C\Gamma_\mu)_{\beta\gamma}}{2z_{14}z_{42}z_{43}z_{23}} \tag{5}$$

where C is the charge conjugation operator.

The correlation function of the bosonic fields gives

$$< i\partial_z X^\mu e^{ik_1\cdot X}(z_1,\overline{z}_1) i\partial_z X^\nu e^{ik_2\cdot X}(z_2,\overline{z}_2) e^{ik_3\cdot X}(z_3,\overline{z}_3) e^{ik_4\cdot X}(z_4,\overline{z}_4) >= \prod_{i<j} |z_{ij}|^{\alpha' k_i\cdot k_j} V^{\mu\nu} \tag{6}$$

where

$$V^{\mu\nu} = \frac{\alpha'\eta^{\mu\nu}}{2z_{12}^2} + \frac{\alpha'^2}{4}\sum_{i\neq1,j\neq2}\frac{k_i^\mu k_j^\nu}{z_{1i}z_{2j}} \tag{7}$$

Note that in the three-point function of massless states, one can use transversality ($e\cdot k = 0$) to simplify the expression (see eqn 12.4.12 of [60]). The transversality condition comes from

the physical state conditions: $L_m - a\delta_{m,0} = G_r = 0$ for $m \geq 0$, $r \geq 0$. It is easy to check that the physical state condition also gives $e \cdot k = 0$ for the massive states we consider here.

For the $SO(32)$ degrees of freedom:

$$< e^{i\frac{K_1}{\sqrt{2\alpha'}}\cdot\tilde{X}(\overline{z}_1)} e^{i\frac{K_2}{\sqrt{2\alpha'}}\cdot\tilde{X}(\overline{z}_2)} e^{i\frac{K_3}{\sqrt{2\alpha'}}\cdot\tilde{X}(\overline{z}_3)} e^{i\frac{K_4}{\sqrt{2\alpha'}}\cdot\tilde{X}(\overline{z}_4)} >= \prod_{i<j} \overline{z}_{ij}^{K_i\cdot K_j} \tag{8}$$

To calculate the four-point amplitude, we can set $z_1 = 0$, $z_2 = x$, $z_3 = 1$, $z_4 = \infty$, and then sum over all permutations. The momentum dependence can be simplified as usual with the Mandelstam s, t, u variables.

Let us focus on one of the amplitudes (the ordering 1234):

$$< c\overline{c}(z_1)c\overline{c}(z_3)c\overline{c}(z_4) >= (z_4\overline{z}_4)^2 \tag{9}$$

$$< \mathcal{V}_\alpha(z_1)\mathcal{V}_\beta(z_2)\mathcal{V}_\gamma(z_3)\mathcal{V}_\delta(z_4) >= \frac{(\overline{u}_1\cdot u_2)(\overline{u}_3\cdot u_4)}{2x(1-x)z_4^2} + \frac{(\overline{u}_1\cdot u_3)(\overline{u}_4\cdot u_2)}{2(1-x)z_4^2} \tag{10}$$

$$V^{\mu\nu} = \frac{\alpha'\eta^{\mu\nu}}{2x^2} + \frac{\alpha'}{4}\left(-\frac{k_2^\mu k_1^\nu}{x^2} - \frac{k_3^\mu k_1^\nu}{x} + \frac{k_2^\mu k_3^\nu}{x(1-x)} + \frac{k_3^\mu k_3^\nu}{1-x}\right) \tag{11}$$

$$\begin{aligned} \prod_{i<j} |z_{ij}|^{\alpha' k_i\cdot k_j} &= |z_4|^{-\alpha' k_4^2}(x\overline{x})^{\alpha' k_1\cdot k_2/2}\left[(1-x)(1-\overline{x})\right]^{\alpha' k_2\cdot k_3/2} \\ &= (x\overline{x})^{-\alpha' s/4}\left[(1-x)(1-\overline{x})\right]^{-\alpha' u/4} \end{aligned} \tag{12}$$

$$\prod_{i<j} \overline{z}_{ij}^{K_i\cdot K_j} = \overline{z_4}^{-K_4^2}\overline{x}^{K_1\cdot K_2}(1-\overline{x})^{K_2\cdot K_3} = \overline{z}_4^{-2}\overline{x}^{-S/2}(1-\overline{x})^{-U/2} \tag{13}$$

Note that the powers of z_4 and $\overline{z}_4$ are canceled. There are many terms in the final expression. Let's focus on the term proportional to $\eta^{\mu\nu}$:

$$C_0\hat{N}^4(e_1\cdot e_2)\int d^2x x^{-\alpha' s/4-1}\overline{x}^{-\alpha' s/4-S/2-2}(1-x)^{-\alpha' u/4}(1-\overline{x})^{-\alpha' u/4-U/2+1} \tag{14}$$

We can express the amplitude in terms of Γ-functions and calculate the correction to the tree-level field theory result. The integral we have to do is

$$\int_C d^2z z^{-\lambda_s-1}\overline{z}^{-\lambda_s-S/2-2}(1-z)^{-\lambda_u}(1-\overline{z})^{-\lambda_u-U/2+1}, \tag{15}$$

where $\lambda_s = \alpha' s/4$, $\lambda_u = \alpha' u/4$. We will suppress the $SO(32)$ structure. Consider the following choice of internal momenta

$$\begin{aligned} K_1 &= -K_2 = (+\frac{1}{2}, ..., +\frac{1}{2}), \\ K_3 &= -K_4 = (0, ...0, +1, 0, ..., 0, +1, 0, ...). \end{aligned} \tag{16}$$

With this choice, we have $S = 0$, $U = -4$. The integral reduces to

$$\int_C d^2 z z^{-\lambda_s - 1} \bar{z}^{-\lambda_s - 2} (1 - z)^{-\lambda_u} (1 - \bar{z})^{-\lambda_u + 3}.. \tag{17}$$

We can choose the following space-time momenta:

$$\begin{aligned} k_1 &= (E, 0, ...0, k) \\ k_2 &= (E, 0, ...0, -k) \\ k_3 &= (-E, 0, ... - E\cos\theta, -E\sin\theta) \\ k_3 &= (-E, 0, ...E\cos\theta, E\sin\theta). \end{aligned} \tag{18}$$

For example, at fixed angle $\theta = \pi/2$, we have

$$\begin{aligned} s &= 4E^2 \\ u &= -2E^2 + m^2, \end{aligned} \tag{19}$$

where $m^2 = 4/\alpha'$.

The $s = 0$ pole could be extracted as follows. First, we expand the integrand of Eq. (17) around $z = 0$ and obtain

$$\begin{aligned} I \sim \int_C d^2 z \ &\Big[(z\bar{z})^{-\lambda_s - 1} \bar{z}^{-1} + (z\bar{z})^{-\lambda_s - 1} \bar{z}^{-1} z + ... \\ &+ (\lambda_u - 3)(z\bar{z})^{-\lambda_s - 1} + ... \Big] \end{aligned} \tag{20}$$

The terms such as the ones displayed on the first line do not contribute to the pole. Therefore, the $s = 0$ pole is from the second line:

$$-\frac{1}{2}(\lambda_u - 3)\frac{1}{\lambda_s}. \tag{21}$$

The full amplitude is

$$\mathcal{A} \propto 2\pi \times \frac{\Gamma(\lambda_s + \lambda_u - 2)\Gamma(-\lambda_s)\Gamma(1 - \lambda_u)}{\Gamma(\lambda_u - 3)\Gamma(\lambda_s + 2)\Gamma(1 - \lambda_s - \lambda_u)}. \tag{22}$$

We can use the Γ-function identity $\Gamma(x)\Gamma(1 - x) = \pi/\sin(\pi x)$ to rewrite this amplitude as

$$\begin{aligned} \mathcal{A} \propto \ & 2 \times \frac{\Gamma(\lambda_s + \lambda_u - 2)\Gamma(1 - \lambda_u)\Gamma(4 - \lambda_u)\Gamma(\lambda_s + \lambda_u)}{\Gamma(\lambda_s + 2)\Gamma(\lambda_s)} \\ \times \ & \frac{\sin[\pi(4 - \lambda_u)]\sin[\pi(\lambda_s + \lambda_u)]}{\sin[\pi\lambda_s]} \times \left(\frac{-1}{\lambda_s}\right). \end{aligned} \tag{23}$$

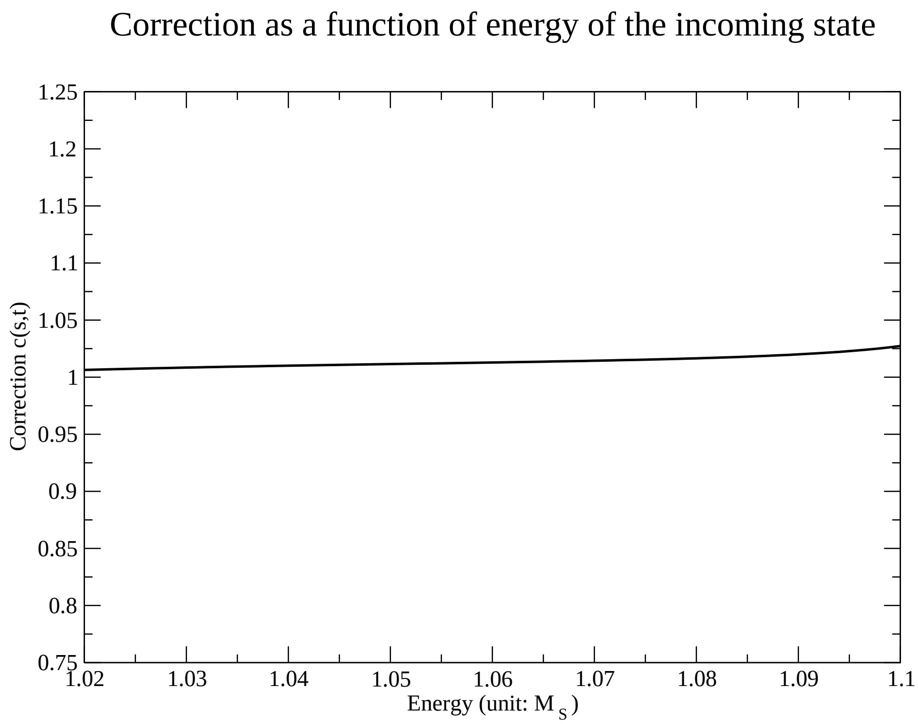

FIG. 5: The corrections of the form factor.

From this we clearly see the closed string poles at $\lambda_s = 4, 5, 6.....$, which correspond to the production of on-shell states in this channel.

We could write the amplitude in the form of $\mathcal{A} = f(u,t)/s + g(s,t,u,\alpha')$, where $g(s,t,u,\alpha')$ is the stringy correction to the zero mass gauge boson exchange in field theory. We could define a measure of the stringy corrections as follows:

$$\frac{g(s,t,u,\alpha')}{f(u,t)/s} = \frac{\mathcal{A}}{f(u,t)/s} - 1, \tag{24}$$

This quantity as a function of energy E is shown in Fig. 5 from which we see that the correction is of order one in a wide range of energies.

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[1] P. Horava and E. Witten, Nucl. Phys. B **460**, 506 (1996) [arXiv:hep-th/9510209]; Nucl. Phys. B **475**, 94 (1996) [arXiv:hep-th/9603142].

[2] N. Arkani-Hamed, S. Dimopoulos, G. R. Dvali, Phys. Lett. B **429**, 263 (1998) [arXiv:hep-ph/9803315]; Phys. Rev. D **59**, 086004 (1999) [arXiv:hep-ph/9807344]; I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, G. R. Dvali, Phys. Lett. B **436**, 257 (1998) [arXiv:hep-ph/9804398].

[3] G. Shiu and S.-H.H. Tye, Phys. Rev. **D58** (1998) 106007 [arXiv:hep-th/9805157].

[4] J. D. Lykken, Phys. Rev. D **54**, 3693 (1996) [arXiv:hep-th/9603133].

[5] A. Lukas, B. A. Ovrut, K. S. Stelle, D. Waldram, Phys. Rev. D **59**, 086001 (1999) [arXiv:hep-th/9803235].

[6] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999) [arXiv:hep-ph/9905221]; Phys. Rev. Lett. **83**, 4690 (1999) [arXiv:hep-th/9906064].

[7] R. Minasian and G. W. Moore, JHEP **9711**, 002 (1997) [arXiv:hep-th/9710230].

[8] E. Witten, JHEP **9812**, 019 (1998) [arXiv:hep-th/9810188].

[9] P. Horava, Adv. Theor. Math. Phys. **2**, 1373 (1999) [arXiv:hep-th/9812135].

[10] See, e.g., E. Witten, arXiv:hep-th/0212247, and references therein.

[11] M. R. Douglas, JHEP **0305**, 046 (2003) [arXiv:hep-th/0303194].

[12] A. Sen, JHEP **9806**, 007 (1998) [arXiv:hep-th/9803194]; JHEP **9808**, 010 (1998) [arXiv:hep-th/9805019]; JHEP **9809**, 023 (1998) [arXiv:hep-th/9808141]; JHEP **9812**, 021 (1998) [arXiv:hep-th/9812031].

[13] O. Bergman and M. R. Gaberdiel, Phys. Lett. B **441**, 133 (1998) [arXiv:hep-th/9806155].

[14] M. Frau, L. Gallot, A. Lerda and P. Strigazzi, Nucl. Phys. B **564**, 60 (2000) [arXiv:hep-th/9903123].

[15] M. R. Gaberdiel and B. J. Stefanski, Nucl. Phys. B **578**, 58 (2000) [arXiv:hep-th/9910109].

[16] B. Stefanski, Nucl. Phys. B **589**, 292 (2000) [arXiv:hep-th/0005153].

[17] N. Quiroz and B. Stefanski, Phys. Rev. D **66**, 026002 (2002) [arXiv:hep-th/0110041].

[18] V. Braun and B. Stefanski, arXiv:hep-th/0206158.

[19] A. Sen, JHEP **9810**, 021 (1998) [arXiv:hep-th/9809111].

[20] L. Gallot, A. Lerda and P. Strigazzi, Nucl. Phys. B **586**, 206 (2000) [arXiv:hep-th/0001049].

[21] M. Cvetic, G. Shiu and A. M. Uranga, Phys. Rev. Lett. **87**, 201801 (2001) [arXiv:hep-th/0107143].

[22] M. Cvetic, G. Shiu and A. M. Uranga, Nucl. Phys. B **615**, 3 (2001) [arXiv:hep-th/0107166].

[23] M. Cvetic, G. Shiu and A. M. Uranga, arXiv:hep-th/0111179.

[24] M. Cvetic, I. Papadimitriou and G. Shiu, arXiv:hep-th/0212177.

[25] M. Berkooz, M. R. Douglas and R. G. Leigh, Nucl. Phys. B **480**, 265 (1996) [arXiv:hep-th/9606139].

[26] R. Blumenhagen, L. Goerlich, B. Kors and D. Lust, JHEP **0010**, 006 (2000) [arXiv:hep-th/0007024].

[27] C. Angelantonj and A. Sagnotti, arXiv:hep-th/0010279.

[28] G. Aldazabal, S. Franco, L. E. Ibanez, R. Rabadan and A. M. Uranga, JHEP **0102**, 047 (2001) [arXiv:hep-ph/0011132]; J. Math. Phys. **42**, 3103 (2001) [arXiv:hep-th/0011073].

[29] R. Blumenhagen, B. Kors and D. Lust, JHEP **0102**, 030 (2001) [arXiv:hep-th/0012156].

[30] L. E. Ibanez, F. Marchesano and R. Rabadan, JHEP **0111**, 002 (2001) [arXiv:hep-th/0105155].

[31] R. Blumenhagen, B. Kors, D. Lust and T. Ott, Nucl. Phys. B **616**, 3 (2001) [arXiv:hep-th/0107138].

[32] S. Forste, G. Honecker and R. Schreyer, Nucl. Phys. B **593**, 127 (2001) [arXiv:hep-th/0008250].

[33] C. Bachas, arXiv:hep-th/9503030.

[34] C. Angelantonj, I. Antoniadis, E. Dudas and A. Sagnotti, Phys. Lett. B **489**, 223 (2000) [arXiv:hep-th/0007090].

[35] J. Polchinski and E. Witten, Nucl. Phys. B **460**, 525 (1996) [arXiv:hep-th/9510169].

[36] S. H. Shenker,arXiv:hep-th/9509132.

[37] M. R. Douglas, D. Kabat, P. Pouliot and S. H. Shenker,Nucl. Phys. B **485**, 85(1997)[arXiv:hep-th/9608024].

[38] S. Chang, C. Coriano and A. E. Faraggi, Nucl. Phys. B **477**, 65 (1996) [arXiv:hep-ph/9605325].

[39] D. J. Chung, E. W. Kolb and A. Riotto, Phys. Rev. D **59**, 023501 (1999) [arXiv:hep-ph/9802238].

[40] D. J. Chung, E. W. Kolb and A. Riotto, Phys. Rev. Lett. **81**, 4048 (1998) [arXiv:hep-ph/9805473].

[41] D. J. Chung, arXiv:hep-ph/9808323.

[42] D. J. Chung, E. W. Kolb and A. Riotto, Phys. Rev. D **60**, 063504 (1999) [arXiv:hep-ph/9809453].

[43] D. J. Chung, arXiv:hep-ph/9809489.

[44] E. W. Kolb, D. J. Chung and A. Riotto, arXiv:hep-ph/9810361.

[45] AGASA homepage, http://www.icrr.u-tokyo.ac.jp/as/as.html

[46] HiRes homepage, http://www.cosmic-ray.org/

[47] K. Greisen, Phys. Rev. Lett. **16**, 748 (1966).

[48] G. T. Zatsepin and V. A. Kuzmin, JETP Lett. **4**, 78 (1966) [Pisma Zh. Eksp. Teor. Fiz. **4**, 114 (1966)].

[49] M. Birkel and S. Sarkar, Astropart. Phys. **9**, 297 (1998) [arXiv:hep-ph/9804285], S. Sarkar and R. Toldra, Nucl. Phys. B **621**, 495 (2002) [arXiv:hep-ph/0108098].

[50] For recent reviews, see S. Sarkar, arXiv:hep-ph/0005256, S. Sarkar, arXiv:hep-ph/0202013.

[51] C. Barbot, M. Drees, F. Halzen and D. Hooper, Phys. Lett. B **555**, 22 (2003) [arXiv:hep-ph/0205230]. C. Barbot, M. Drees, F. Halzen and D. Hooper, Phys. Lett. B **563**, 132 (2003) [arXiv:hep-ph/0207133].

[52] For a review, see K. Griest and M. Kamionkowski, Phys. Rept. **333** (2000) 167.

[53] R. Rabadan and G. Shiu, JHEP **0305**, 045 (2003) [arXiv:hep-th/0212144].

[54] G. R. Dvali and S. H. Tye, Phys. Lett. B **450**, 72 (1999) [arXiv:hep-ph/9812483].

[55] N. Jones, H. Stoica and S. H. Tye, JHEP **0207**, 051 (2002) [arXiv:hep-th/0203163]; Phys. Lett. B **563**, 6 (2003) [arXiv:hep-th/0303269].

[56] S. Sarangi and S. H. Tye, Phys. Lett. B **536**, 185 (2002) [arXiv:hep-th/0204074].

[57] G. Shiu, S. H. Tye and I. Wasserman, Phys. Rev. D **67**, 083517 (2003) [arXiv:hep-th/0207119].

[58] G. Shiu, arXiv:hep-th/0210313.

[59] L. Pogosian, S. H. Tye, I. Wasserman and M. Wyman, Phys. Rev. D **68**, 023506 (2003) [arXiv:hep-th/0304188].

[60] J. Polchinski, *String Theory*, Cambridge University Press (1998).

[61] Magnetic monopoles are common in string models and their cosmological implications have been explored [10]. As we will discuss, there are similarities and differences between D-matter and magnetic monopoles.

[62] The disk diagram has a normalization of g_s^{-1}. Each open string vertex operator insertion carries a factor of $g_s^{1/2}$. Hence the overall amplitude of such diagram is proportional to $g_s^{1/2}$.

[63] Note however that some D-matter states could act as magnetic point sources of the worldvolume gauge field on some higher-dimensional branes, in which case their properties are similar to that of 't Hooft Polyakov magnetic monopoles. For example, a D1-brane ending on two D3-branes corresponds to a magnetic monopole of the $SU(2)$ gauge theory on the D3-brane worldvolume. This configuration is S-dual to a fundamental string ending on the D3-branes. The power counting of g_s for D-matter interactions can be obtained by taking $g_s \to 1/g_s$ of the S-dual amplitudes.

[64] D-matter states could be bosonic or fermionic. They correspond to the bosonic and fermionic zero modes of the stable D-branes respectively.

[65] Notice that charge conservation with respect to the unbroken gauge group on the worldvolume of D-matter has to be satisfied if only one endpoint of the open strings is attached to it. In principle, the gauge symmetry on the world volume of D-matter could be broken in a realistic model so this type of requirement would not impose a restriction on the possible types of excited states

[66] At a more technical level, the zeroth mode of the Virasoro generator of an excited D-particle state is $L_0 = L_0^D + L_0^\mathcal{V}$. Requiring L_0 annihilate the D-matter state will give us the mass formula.

[67] Whether the open string states carry momentum/winding numbers depends on whether the corresponding compact dimension is along or transverse to the worldvolumes of D-matter and the Standard Model (SM) branes. If a compact dimension is along the worldvolumes of both the D-matter and the SM branes (Neumann-Neumann boundary condition), the open string states can carry momentum but not winding numbers. Conversely, if a compact dimension is transverse to both the D-matter and the SM branes (Dirichlet-Dirichlet boundary condition), the open string states can carry winding but not momentum numbers. Finally, if a compact dimension is transverse to either the D-matter or the SM branes but not both (Dirichlet-Neumann boundary condition), the open string states do not carry any momentum or winding numbers.